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Weak Scale Superstrings

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Recent developments in string duality suggest that the string scale may not be irrevocably tied to the Planck scale. Two explicit but unrealistic examples are described where the ratio of the string scale to the Planck scale is arbitrarily small. Solutions which are more realistic may exist in the intermediate coupling or "truly strong coupling" region of the heterotic string. Weak scale superstrings have dramatic experimental consequences for both collider physics and cosmology.

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1. Introduction

The discovery of string dualities is reshaping the way we think about string theory. Indeed even the terminology "string theory" has become suspect, given the apparent dualities between certain string compactifications and compactifications of eleven-dimensional "M-theory" [1,2,3,4] or twelve-dimensional "F-theory" [5,6]. The heterotic, Type II, Type I, and Type I′ superstrings are dual descriptions of the same underlying theory.

In light of these radical developments it is important to reexamine our understanding of how string theory is likely to be related to the real world. A step in this direction is the recent paper by Witten [7]. He observes that superstring phenomenology to date has assumed certain relationships between parameters which hold in the weak coupling regime of the heterotic string, but which may *not* be valid generally. In particular there is the famous tree-level formula [8]

$$\alpha' M_{\rm P}^2 = \frac{4}{k\alpha_{\rm U}} \quad . \tag{1.1}$$

Here α' is the string tension (which has units of length squared); for simplicity we will define $m_s = 1/\sqrt{\alpha'}$ to be the string scale. $M_{\rm P}$ is the Planck mass $\simeq 10^{19}$ GeV defined from Newton's constant by $G_{\rm N} = 1/M_{\rm P}^2$. $\alpha_{\rm U} = g_{\rm U}^2/4\pi$, where $g_{\rm U}$ is the unified gauge coupling. The parameter k is the Kac-Moody level; it is compactification dependent but of order one [9]. If the group is nonsimple k takes independent values for each group factor.

Since the value of $g_{\rm U}$ is presumably of order one, this implies that the string scale m_s is not far below the Planck scale. The string scale determines both the scale of gauge coupling unification and the scale of Regge recurrences (the massive string modes). These are thus both predicted to be in the range 10^{17} – 10^{18} GeV.

Reference [7] points out that this relationship of scales and couplings can be radically altered in the strong coupling regime of the heterotic string. This is shown by a duality map of the strong coupling $SO(32)$ or $E_8 \times E_8$ heterotic strings, compactified to four dimensions, to (respectively) a weak coupling Type I string compactified to four dimensions, or M-theory compactified first to $R^{10} \times S^1/Z_2$, then to four dimensions. For the $SO(32)$ string one finds that

$$m_s^2/M_{\rm P}^2 \propto \lambda_I \tag{1.2}$$

where λ_I is the ten-dimensional Type I string coupling, determined dynamically by the vacuum expectation value of the dilaton. Since we are in the weak coupling regime for the Type I string (1.2) can imply small values of $m_s^2/M_{\rm P}^2$.

In the $E_8\times E_8$ case one finds that

$$m_s^2/M_{\rm P}^2 \propto \kappa^{2/9}/\rho \tag{1.3}$$

where κ is the eleven-dimensional gravitational coupling and ρ is the compactification radius in $R^{10}\times S^1$. Here the story is more complicated, but in [7] it is shown that, for the symmetric embedding of the gauge bundle, the ratio $m_s^2/M_{\rm P}^2$ can also be small consistent with the assumption that the ten-dimensional fields are weakly coupled.

If the string scale is not irrevocably tied to the Planck scale, it is natural to explore the idea that it may instead be tied to the electroweak scale (246 GeV). I will use the name *weak scale superstrings* to denote string solutions with m_s in the range from 250 GeV up to a few TeV.

2. An Example in Six Dimensions

Weak scale superstrings are a subset of the class of string solutions for which the ratio $m_s/M_{\rm P}$ can be tuned arbitrarily small while keeping (at least some) gauge couplings of order one. In six dimensions the gauge coupling has dimensions of length; this defines an energy scale below which the six-dimensional effective gauge theory is weakly coupled. Thus one can examine the six-dimensional analog of weak scale superstrings by looking for solutions where

$$\frac{(\alpha')^2}{\kappa^2} \gg 1; \qquad \frac{\alpha'}{g^2} \sim {\rm O}(1) \quad . \tag{2.1}$$

where κ is the six-dimensional gravitational coupling.

There are two reasons for considering six-dimensional examples first. One is that, given a six-dimensional solution which satisfies (2.1), we can in general obtain four-dimensional solutions of the type we want by further compactifying two dimensions at a compactification scale which is of order one in string units. More importantly, in six dimensions the constraints from both anomaly cancellation and N=1 spacetime supersymmetry are more severe than in four dimensions. This allows one to extract information more reliably from the interesting region of moduli space.

The first example I will discuss is a six-dimensional compactification of the Type I superstring on a $K3$ Z_2 orbifold, a class of solutions recently constructed by Gimon and Polchinski [10]. These solutions have N=1 (more precisely, (0,1)) spacetime supersymmetry, the minimal amount of supersymmetry in six dimensions. A toroidal compactification of such a solution to four dimensions will produce solutions with N=2 supersymmetry. In the case where all sixteen of the Dirichlet 5-branes are at a fixed point of the orbifold projection, the gauge group is $U(16)\times U(16)$. The first/second $U(16)$ is carried by Chan-Paton factors associated with open strings with ends attached to Dirichlet 9-branes/5-branes, respectively. Moving all sixteen 5-branes away from the fixed point and turning on appropriate Wilson lines gives a very similar solution with gauge group $USp(16)\times USp(16)$ [11].

The massless particle content consists of the gravity multiplet, one tensor multiplet, 20 gauge singlet hypermultiplets, the vector multiplets of $U(16)\times U(16)$, and hypermultiplets transforming under $U(16)\times U(16)$ as a $(16,16)$, a $(120+\overline{120},1)$, and a $(1,120+\overline{120})$.

Anomaly cancellation and spacetime supersymmetry fix completely the form of certain terms in the effective low energy field theory action [12,13,4]. Thus in the Einstein frame the action is:

$$\begin{aligned}\frac{(2\pi)^3}{(\alpha')^2}\int d^6x\,\sqrt{g}\Big\{\, R-\frac{1}{12}\mathrm{e}^{-2\phi}H^2 \\ -\frac{\alpha'}{8}\sum_{\alpha=1,2}\left(v_\alpha \mathrm{e}^{-\phi}+\tilde{v}_\alpha \mathrm{e}^{\phi}\right)\mathrm{tr}\,F_\alpha^2+\ldots\Big\}\quad .\end{aligned} \tag{2.2}$$

Here ϕ is the scalar component of the tensor multiplet, R is the Ricci scalar, H is the 3-form field strength, and F_1, F_2 are the $U(16)\times U(16)$ field strengths.

Furthermore, the parameters v_1, v_2, $\tilde{v}_1$, $\tilde{v}_2$ are fixed by anomaly cancellation*. The anomaly 8-form can be written [15]:

$$I_8=(\mathrm{tr}\,R^2)^2+\frac{1}{6}\mathrm{tr}\,R^2\sum_\alpha X_\alpha^{(2)}-\frac{2}{3}\sum_\alpha X_\alpha^{(4)}+4\sum_{\alpha<\beta}Y_{\alpha\beta}, \tag{2.3}$$

where

$$\begin{aligned}X_\alpha^{(n)}&=\mathrm{Tr}\,F_\alpha^n-\sum_i n_i \mathrm{tr}_i\,F_\alpha^n \\ Y_{\alpha\beta}&=\sum_{ij}n_{ij}\,\mathrm{tr}_i\,F_\alpha^2\,\mathrm{tr}_j\,F_\beta^2.\end{aligned} \tag{2.4}$$

* For simplicity I will ignore the U(1) anomalies. For a complete analysis, see [14].

Here the symbol Tr denotes a trace in the adjoint representation and tr_i denotes a trace in the representation R_i (of the simple group G_α). n_i is the number of hypermultiplets in the representation R_i of G_α and n_{ij} is the number of representations (R_i, R_j) of $G_\alpha \times G_\beta$ which occur. The Green-Schwarz anomaly cancellation mechanism requires that the anomaly 8-form should factorize as

$$I_8 = (\mathrm{tr}\, R^2 - \sum_\alpha v_\alpha \mathrm{tr}\, F_\alpha^2)(\mathrm{tr}\, R^2 - \sum_\alpha \tilde{v}_\alpha \mathrm{tr}\, F_\alpha^2) \quad , \tag{2.5}$$

where tr denotes the trace in the fundamental representation.

Using the trace identities of ref. [16], one finds for the $U(16)\times U(16)$ model

$$\begin{aligned} X_1^{(2)} &= -12\mathrm{tr}\, F_1^2 \quad , \\ X_2^{(2)} &= -12\mathrm{tr}\, F_2^2 \quad , \\ X_1^{(4)} &= X_2^{(4)} = 0 \quad , \\ Y &= \mathrm{tr}\, F_1^2 \mathrm{tr}\, F_2^2 \quad . \end{aligned} \tag{2.6}$$

Thus

$$I_8 = \left(\mathrm{tr}\, R^2 - 2\mathrm{tr}\, F_1^2\right)\left(\mathrm{tr}\, R^2 - 2\mathrm{tr}\, F_2^2\right) \tag{2.7}$$

which implies:

$$\begin{aligned} v_1 &= 2\,, \qquad \tilde{v}_1 = 0 \\ v_2 &= 0\,, \qquad \tilde{v}_2 = 2 \end{aligned} \tag{2.8}$$

The result v_2=0 indicates that the gauge bosons of the second $U(16)$ are inherently non-perturbative. This is expected as they are associated with the Dirichlet 5-branes [4,17].

Let us now rescale from the Einstein frame to the string metric frame; this is the frame in which m_s actually sets the scale of the Regge recurrences. Rescale the metric by

$$g_{\mu\nu} \to \frac{\mathrm{e}^{-\phi}}{\lambda_I} g_{\mu\nu} \tag{2.9}$$

where λ_I is the ten-dimensional Type I string coupling. Then (2.2) becomes:

$$\begin{aligned} \frac{(2\pi)^3}{(\alpha')^4} \int d^6x \sqrt{g}\, V_I \Big\{ & \frac{1}{\lambda_I^2} R - \frac{1}{12} \mathrm{e}^{-2\phi} H^2 \\ & - \frac{\alpha'}{4\lambda_I} \mathrm{tr}\, F_1^2 - \frac{(\alpha')^3}{4\lambda_I V_I} \mathrm{tr}\, F_2^2 + \ldots \Big\} \quad . \end{aligned} \tag{2.10}$$

where

$$V_I \equiv \mathrm{e}^{-2\phi} (\alpha')^2 \tag{2.11}$$

can be regarded as the effective compactification volume; note this analysis in no way depends on an implicit assumption that V_I is large.

From (2.10) we can read off the six-dimensional gravitational and gauge couplings:

$$\begin{aligned}\frac{(\alpha')^2}{\kappa^2} &\sim \frac{V_I}{\lambda_I^2(\alpha')^2} \\ \frac{\alpha'}{g_1^2} &\sim \frac{V_I}{\lambda_I(\alpha')^2} \\ \frac{\alpha'}{g_2^2} &\sim \frac{1}{\lambda_I}\end{aligned} \tag{2.12}$$

The analog of weak scale superstrings thus corresponds to very weak coupling and small V_I:

$$\lambda_I \ll 1\,, \qquad V_I/(\alpha')^2 = \mathrm{O}(\lambda_I) \tag{2.13}$$

In this region of moduli space we then have:

$$\frac{(\alpha')^2}{\kappa^2} \sim \frac{1}{\lambda_I}\,, \quad \frac{\alpha'}{g_1^2} \sim \mathrm{O}(1)\,, \quad \frac{\alpha'}{g_2^2} \sim \frac{1}{\lambda_I}\,. \tag{2.14}$$

There are two widely separated energy scales. The lower scale is the scale at which the first Regge recurrences appear and at which the first $U(16)$ gauge coupling gets strong. The higher scale is the scale at which both gravity and the second $U(16)$ gauge coupling get strong. In this analogy the standard model gauge group would be embedded in the first $U(16)$.

It is also instructive to look at the equivalent heterotic or Type I′ description of these solutions. The table below show how the string couplings and compactification scales are related by duality [18,19]:

Heterotic	Type I	Type I′	
$\frac{1}{\lambda_h}$	λ_I	$\frac{(\alpha')^2\lambda_{I'}}{V_{I'}}$	(2.15)
$\lambda_I^2 V_h$	V_I	$\frac{(\alpha')^4}{V_{I'}}$	

In the heterotic description, we are in a region of strong coupling and large radius. In the Type I′ description we are also at large radius, but the ten-dimensional string coupling is of order one.

3. Another Example in Six Dimensions

Another simple example comes from the $SO(32)$ heterotic string compactified on $K3$. The $K3$ compactification requires a gauge bundle with instanton number 24. As shown by Witten [3], at the special region in moduli space where all 24 instantons shrink to zero size, the gauge group is enhanced to $SO(32)\times Sp(24)$. The extra $Sp(24)$ gauge bosons are inherently nonperturbative and are associated with solitonic 5-branes, just as the second $U(16)$ in the $K3$ orbifold discussed above was associated with the dual Dirichlet 5-branes.

Because of anomaly cancellation and supersymmetry the low energy effective action in the Einstein frame has the same form as (2.2). The v, $\tilde{v}$ parameters are determined to be [20]:

$$
\begin{aligned}
v_{32} &= 1\,, \qquad & \tilde{v}_{32} &= -2 \\
v_{24} &= 0\,, \qquad & \tilde{v}_{24} &= 2
\end{aligned}
\tag{3.1}
$$

For the heterotic string e^{ϕ} is the six-dimensional effective string coupling, i.e.

$$
\mathrm{e}^{2\phi} = \frac{\lambda_h^2}{V_h}
\tag{3.2}
$$

where V_h is the volume of $K3$ and λ_h is the ten-dimensional string coupling. Thus the proper rescaling from the Einstein frame to the string metric frame is given by:

$$
g_{\mu\nu} \to \mathrm{e}^{-2\phi} g_{\mu\nu}
\tag{3.3}
$$

Then (2.2) becomes:

$$
\begin{aligned}
\frac{(2\pi)^3}{(\alpha')^4} \int d^6x\, \sqrt{g}\, V_h \Big\{ &\frac{1}{\lambda_h^2} R - \frac{1}{12\lambda_h^2} H^2 \\
&- \frac{\alpha'}{8\lambda_h^2}\left(1 - \frac{2(\alpha')^2\lambda_h^2}{V_h}\right) \mathrm{tr}\, F_{32}^2 - \frac{(\alpha')^3}{4V_h} \mathrm{tr}\, F_{24}^2 + \ldots \Big\} \quad .
\end{aligned}
\tag{3.4}
$$

From (3.4) we can read off the six-dimensional gravitational and gauge couplings:

$$
\begin{aligned}
\frac{(\alpha')^2}{\kappa^2} &\sim \frac{V_h}{\lambda_h^2(\alpha')^2} \\
\frac{\alpha'}{g_{32}^2} &\sim \frac{V_h}{\lambda_h^2(\alpha')^2} - 2 \\
\frac{\alpha'}{g_{24}^2} &\sim 1
\end{aligned}
\tag{3.5}
$$

Let us then consider the case where the ten-dimensional string coupling λ_h is of order one, while the heterotic volume V_h is large. In this region of moduli space we then have:

$$\frac{(\alpha')^2}{\kappa^2} \sim \frac{V_h}{(\alpha')^2}\,, \quad \frac{\alpha'}{g_{32}^2} \sim \frac{V_h}{(\alpha')^2}\,, \quad \frac{\alpha'}{g_{24}^2} \sim 1\,. \tag{3.6}$$

There are two widely separated energy scales. The lower scale is the scale at which the first Regge recurrences appear and at which the $Sp(24)$ gauge coupling gets strong. The higher scale is the scale at which both gravity and the $SO(32)$ gauge coupling get strong. In this analogy the standard model gauge group would be embedded in $Sp(24)$.

4. Realistic Weak Scale Superstrings

The six-dimensional examples considered above are very far from a solution which could correspond to a realistic weak scale superstring. One obvious difficulty is that taking the compactification volume to be very large in string units (as in the second example or in the Type I′ picture of the first example) is a phenomenological disaster if the string scale itself is only a TeV. Thus for a realistic model we must suppose that the small ratio m_s/M_P is associated with some modulus which can get a very large vev without generating unwanted observable light states.

Another obvious difficulty is the notorious problem of stablizing the vev of the dilaton [21]. In any weak coupling limit of the superstring, the dilaton vev vanishes – another phenomenological disaster. As discussed by Dine and Shirman [22], a realistic superstring probably must reside in a region of moduli space which admits no weak coupling description. Both six-dimensional examples fail this criterion, the first in the Type I description and the second in the Type I′ description. However this failure is not as bad as it could have been, since in both cases the weak coupling, small radius description is only accessible due to extra symmetry of the compactifications. In a realistic solution we should at any rate avoid extra symmetries which can prevent a stable nonzero dilaton vev even at intermediate and strong coupling [23].

Dine and Shirman [22] have identified a possibly unique region of the moduli space of four-dimensional compactifications which satisfies their criterion without requiring all moduli to take intermediate values. This "truly strong coupling" region corresponds to $\lambda_h \gg 1$, with the 6-dimensional compactification volume V_h scaling like λ_h^2. Thus in the heterotic, Type I, and Type I′ pictures we have:

$$
\begin{array}{cccc}
\text{Heterotic} & \text{Type I} & \text{Type I}' & \\
\lambda_h & \lambda_I = \dfrac{1}{\lambda_h} & \lambda_{I'} = \mathrm{O}(1) & (4.1) \\
V_h \sim \lambda_h^2 & V_I = \mathrm{O}(\lambda_I) & V_{I'} = \mathrm{O}(1/\lambda_I) &
\end{array}
$$

Consider such solutions in the Type I description. If the 6-volume V_I were large instead of small, we would be justified in writing the effective action as:

$$\frac{(2\pi)^3}{(\alpha')^4}\int d^4x\,\sqrt{g}\,V_I\Big\{\,\frac{1}{\lambda_I^2}R - \frac{k\alpha'}{4\lambda_I}\mathrm{tr}\,F^2 + \dots\Big\} \quad . \tag{4.2}$$

where k=1 for large volume. For weak coupling and large volume, we can read off the gauge and gravitational couplings from the tree-level terms in (4.2). As V_I shrinks, this is no longer true, in general. In fact for $V_I{\sim}\lambda_I{\ll}1$, one should regard V_I as representing the scaling of some moduli, but not as a classical volume.

However there is likely to be a large subclass of solutions in the "truly strong coupling" region where the gauge and gravitational couplings are still determined by an effective action of the form (4.2), where V_I is to be regarded as some scaling function of moduli and the parameter k (also a function of some moduli) is of order one. As discussed above, we also must require that the modulus vev that makes V_I small must somehow not also lead to unwanted observable light states. Whether this is likely –or even possible– I do not know.

For these solutions we will have α_{U} of order one while

$$\alpha' M_{\mathrm{P}}^2 = \frac{4}{k\lambda_I\alpha_{\mathrm{U}}} \tag{4.3}$$

i.e. the string scale is arbitrarily smaller than the Planck scale. Thus the "truly strong coupling" region (broadly interpreted) may be a likely place to find realistic weak scale superstrings, if they exist.

5. Objections to Weak Scale Superstrings

5.1. Gauge coupling unification

In reference [7] the results summarized in the introduction were obtained in the context of obtaining a modest reduction in the ratio m_s^2/M_{P}^2 beyond what is implied by (1.1). The motivation is the well known apparent gauge coupling unification at $\sim 10^{16}$ GeV implied by a naive renormalization group evolution of the measured low energy couplings plus minimal SUSY thresholds.

However it is not at all obvious that gauge coupling unification in the string sense has any direct relation to this apparent unification of the standard model gauge couplings. Even for the heterotic string at weak coupling, we know (see the discussion below (1.1)) that gauge coupling unification in the string sense does not necessarily imply *equality* of the gauge couplings at some scale. Thus the only argument pinning the string scale to 10^{16} GeV is the conviction that the apparent unification at that scale is "too close" to be a coincidence. This argument is even weaker than it seems, since it is possible that, while *not* a coincidence, the apparent unification maps into some sophisticated structure of the underlying string theory, without requiring an *actual* field theory desert between 10^3 and 10^{16} GeV.

5.2. The success of weak coupling heterotic models

A number of weak coupling heterotic string models have been built which exhibit an elegant confluence of favorable phenomenological attributes. These models have three generations of standard model chiral fermions, embed the standard model gauge group, and have a natural hidden sector suitable for dynamical supersymmetry breaking. They also exhibit new symmetries which naturally give a hierarchical structure to the Yukawa matrices. For recent reviews, see [24,25,26,27,28].

Thus one could argue that the hypothesis of weak scale superstrings moves us very far away from a class of string solutions which look very much like the real world.

One problem with this argument is that it includes a number of theoretical assumptions in its definition of "the real world". Another problem is that we are only just beginning to understand the principles which control the relationships between phenomenological attributes in such solutions. Some features of these solutions, such as symmetries of the superpotential which restrict Yukawa couplings, should survive if we deform the solutions into the intermediate coupling region [23], where (we hope) the dilaton vev is stabilized. But beyond this it is still premature to use these weak coupling solutions as a way of constraining properties of a realistic string solution.

5.3. Spacetime supersymmetry

Spacetime supersymmetry is motivated in particle theory as a way to stabilize the hierarchy between the electroweak scale and the Planck scale. With superstrings, this ties in nicely with the fact that spacetime SUSY also removes tachyons from the physical string spectrum, and guarantees a vanishing cosmological constant.

If the string scale is around a TeV we lose the original motivation for spacetime SUSY. In fact spacetime supersymmetry becomes a serious problem, since it is notoriously difficult to break supersymmetry in a phenomenologically acceptable way at such a low scale. Furthermore the supersymmetry mass splittings would now be the same order of magnitude as the spacing of the Regge recurrences.

This suggests that a viable weak scale superstring solution may *not* exhibit spacetime supersymmetry in the effective field theory below the string scale.

5.4. Why the electroweak scale?

Why the Planck scale? String dynamics softens the ultraviolet behavior of quantum gravity. With the possible exception of cosmology, I know of no consideration which says that these stringy effects cannot set in at a scale where gravitational forces are still weak. Of course, since the low energy effective field theory action will contain an infinite number of higher dimension terms suppressed by powers of the string scale, weak scale superstrings are constrained somewhat by low energy data -e.g. flavor changing neutral currents. But these constraints are no more severe than for other new physics scenarios at the TeV scale.

6. Experimental Consequences

The hypothesis of weak scale superstrings has spectacular consequences for collider physics at TeV energies. Each of the known particles of the standard model (as well as the graviton) sits at the base of a Regge trajectory. There are an infinite number of Regge recurrences, with progressively higher masses and spins. These particles carry standard model quantum numbers including color and are unstable. The lightest ones could have masses as low as a few hundred GeV without violating current experimental bounds.

An obvious guess for the lightest Regge recurrences are the heavy spin 3/2 partners of light quarks and leptons. For masses in the range from a few hundred GeV to a TeV the heavy spin 3/2 quarks will be easier to detect than the heavy leptons.

The relatively light Regge recurrences may also be accompanied by relatively light Kaluza-Klein modes, if one or more of the effective compactification radii is of order the weak scale rather than the Planck scale. In this case [29] a plausible guess for the lightest Kaluza-Klein modes are the heavy partners of the gluons.

In this regard it is interesting to note that either spin 3/2 heavy quarks or heavy color octets are possible explanations [30,31] of the excess in jet production for $E_t > 200$ GeV reported by the CDF collaboration in $p\bar{p}$ collisions at the Tevatron [32].

The effects of Regge recurrences on the single jet inclusive cross section will resemble the effects of compositeness: in both cases the amplitude has an s/M enhancement at high E_t. However it should be possible to distinguish the higher spin Regge recurrences by examining the jet angular distributions.

If the real world is a weak scale superstring the LHC will produce unintelligible results when operated at design energy and luminosity. It will be necessary in that case to resort to something like a DiTevatron or TEV33 to have any hope of sorting out the superstring threshold region.

Weak scale superstrings also have profound implications for cosmology and black hole physics. The number of heavy string states increases exponentially with mass; this implies a Hagedorn temperature of a few TeV [33]. The existence of such a Hagedorn transition will require a radical rethinking of inflation, structure formation, and baryogenesis.

It will be difficult to construct realistic weak scale superstring models, even if they exist. But if they are there, we will certainly discover them in high energy colliders.

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